HADRONIC TWO-BODY CHARMLESS B DECAYS

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Implications of recent CLEO measurements of hadronic charmless B decays are discussed

1 Effective Wilson Coefficients

In the absence of first-principles calculations for hadronic matrix elements, it is customary to evaluate the matrix elements under the factorization hypothesis so that $\langle O(\mu) \rangle$ is factorized into the product of two matrix elements of single currents, governed by decay constants and form factors. However, the naive factorized amplitude is not renormalization scale- and γ_5 scheme- independent as the scale and scheme dependence of Wilson coefficients are not compensated by that of the factorized hadronic matrix elements. In principle, the scale and scheme problems with naive factorization will not occur in the full amplitude since $\langle O(\mu) \rangle$ involves vertex-type and penguin-type corrections to the hadronic matrix elements of the 4-quark operator renormalized at the scale μ . Formally, one can write

$$\langle O(\mu) \rangle = g(\mu, \mu_f) \langle O(\mu_f) \rangle,$$
 (1)

where μ_f is a factorization scale, and $g(\mu, \mu_f)$ is an evolution factor running from the scale μ to μ_f which is calculable because the infrared structure of the amplitude, if any, is absorbed into $\langle O(\mu_f) \rangle$. Writing

$$c^{\text{eff}}(\mu_f) = c(\mu)q(\mu, \mu_f), \tag{2}$$

the effective Wilson coefficients are formally scheme and μ -scale independent. The factorization approximation is then applied to $\langle O(\mu_f) \rangle$ afterwards.

In principle, one can work with any quark configuration, on-shell or off-shell, to compute the decay amplitude. If the off-shell quark momentum is chosen as the infrared cutoff, $g(\mu, \mu_f)$ will depend on the gauge of the gluon field. This is all right as the gauge dependence belongs to the infrared structure of the wave function so that the physical amplitude is gauge independent ¹. However, if factorization is applied to $\langle O(\mu_f) \rangle$, the information of gauge dependence characterized by the wave function will be lost. Consequently, c^{eff} will depend on the choice of gauge, a difficulty pointed out by Buras and Silvestrini ². As pointed out in ¹, within the factorization framework, one must work in the on-shell quark scheme to obtain gauge invariant and infrared finite

 c_i^{eff} and infrared poles, if any, are absorbed into universal bound-state wave functions. It should be stressed that the constant matrix \hat{r}_V arising from vertex-like corrections is not arbitrary due to the infrared finiteness of vertex-like diagrams: The infrared divergences in individual vertex-type diagrams cancel in their sum. The gauge-invariant \hat{r}_V matrices in naive dimension regularization and 't Hooft-Veltman renormalization schemes are first given in ³ and ⁴ (see also ⁵).

2 Nonfactorized Effects

It is known that the effective Wilson coefficients appear in the factorizable decay amplitudes in the combinations $a_{2i} = c_{2i}^{\text{eff}} + \frac{1}{N_c} c_{2i-1}^{\text{eff}}$ and $a_{2i-1} = c_{2i-1}^{\text{eff}} + \frac{1}{N_c} c_{2i}^{\text{eff}}$ ($i=1,\cdots,5$). Phenomenologically, the number of colors N_c is often treated as a free parameter to model the nonfactorizable contribution to hadronic matrix elements and its value can be extracted from the data of two-body nonleptonic decays. Nonfactorizable effects in the decay amplitudes of $B \to PP$, VP can be absorbed into the parameters a_i^{eff} . This amounts to replacing N_c in a_i^{eff} by $(N_c^{\text{eff}})_i$. Explicitly,

$$a_{2i}^{\text{eff}} = c_{2i}^{\text{eff}} + \frac{1}{(N_c^{\text{eff}})_{2i}} c_{2i-1}^{\text{eff}}, \qquad a_{2i-1}^{\text{eff}} = c_{2i-1}^{\text{eff}} + \frac{1}{(N_c^{\text{eff}})_{2i-1}} c_{2i}^{\text{eff}}, \qquad (3)$$

where $i = 1, \dots, 5$ and

$$(1/N_c^{\text{eff}})_i \equiv (1/N_c) + \chi_i \,,$$
 (4)

with χ_i being the nonfactorizable terms, which receive contributions from nonfactorized vertex-type, penguin-type and spectator corrections. In general, χ_i and $(N_c^{\text{eff}})_i$ are complex.

Naive factorization does not work in the presence of nonfactorized contributions. Nevertheless, if χ_i are universal (i.e. channel by channel independent), then we still have generalized factorization, which is likely to be justified for hadronic charmless B decays due to their large energy release. Since the Fierz transformation of (V-A)(V+A) operators is quite different from that of (V-A)(V-A) operators, we shall assume that $\chi_{LR} \neq \chi_{LL}$, where $\chi_{LL} \equiv \chi_{1,2,3,4,9,10}$ and $\chi_{LR} \equiv \chi_{5,6,7,8}$, or equivalently, $N_c^{\rm eff}(LR) \neq N_c^{\rm eff}(LL)$ with $N_c^{\rm eff}(LL) \equiv \left(N_c^{\rm eff}\right)_{1,2,3,4,9,10}$ and $N_c^{\rm eff}(LR) \equiv \left(N_c^{\rm eff}\right)_{5,6,7,8}$. As shown in $N_c^{\rm eff}(LL) \equiv N_c^{\rm eff}(LL) = N_c^{\rm eff}(LL)$. In principle, $N_c^{\rm eff}(LL) \equiv N_c^{\rm eff}(LL)$ and the theoretical study of nonleptonic rare $N_c^{\rm eff}(LL) \equiv N_c^{\rm eff}(LL)$. In principle, $N_c^{\rm eff}(LL) \equiv N_c^{\rm eff}(LL)$ and the case of charm decay. However, in the energetic two-body $N_c^{\rm eff}(LL) \equiv N_c^{\rm eff}(LL)$ is expected to be process insensitive as supported by the data $N_c^{\rm eff}(LL) \equiv N_c^{\rm eff}(LL)$

The observation $N_c^{\rm eff}(LL) < 3 < N_c^{\rm eff}(LR)$ is theoretically justified by a recent perturbative QCD calculation of charmless B decays in the heavy quark limit. As pointed out in 5 , in the heavy quark limit, all nonfactorizable diagrams are dominated by hard gluon exchange, while soft gluon effects are suppressed by factors of $\Lambda_{\rm QCD}/m_b$. In other words, the nonfactorized term is calculable as expansion in α_s in the heavy quark limit. Following 5 , we find the nonfactorized terms:

$$\chi_{LR} = -\chi_{LL} = -\frac{\alpha_s}{4\pi} \frac{C_F}{N_c} (f^{\rm I} + f^{\rm II}),$$
(5)

where the hard scattering function $f^{\rm I}$ corresponds to hard gluon exchange between the two outgoing light mesons and $f^{\rm II}$ describes the hard nonfactorized effect involving the spectator quark of the B meson. Two remarks are in order. (i) Since $f^{\rm I}$ is complex due to final-state interactions via hard gluon exchange 5 , so are χ_i and $N_c^{\rm eff}(LL)$ and $N_c^{\rm eff}(LR)$. Nevertheless, the complex phases of χ_i are in general small. (ii) Because Re $\chi_{LL}>0$, it is obvious that $N_c^{\rm eff}(LL)<3$ and $N_c^{\rm eff}(LR)>3$. Furthermore, $N_c^{\rm eff}(LL)\sim2$ implies $N_c^{\rm eff}(LR)\sim6$. Therefore, the empirical observation $N_c^{\rm eff}(LR)>3>N_c^{\rm eff}(LL)$ shown in 3 gets a firm justification from the perturbative QCD calculation.

3 Tree-Dominated Charmless B Decays

CLEO has observed several tree-dominated charmless B decays which proceed at the tree level through the b quark decay $b \to u\bar{u}d$ and at the loop level via the $b \to d$ penguin diagrams: $B \to \pi^+\pi^-$, $\rho^0\pi^\pm$, $\omega\pi^\pm$, $\rho^\pm\pi^\mp$. The updated branching ratios have been reported at this Conference ⁶:

$$\mathcal{B}(B^0 \to \pi^+ \pi^-) = (4.3^{+1.6}_{-1.4} \pm 0.6) \times 10^{-6},$$

$$\mathcal{B}(B^\pm \to \rho^0 \pi^\pm) = (10.4^{+3.3}_{-3.4} \pm 2.1) \times 10^{-6},$$

$$\mathcal{B}(B^\pm \to \omega \pi^\pm) = (11.3^{+3.3}_{-2.9} \pm 1.5) \times 10^{-6},$$

$$\mathcal{B}(B^0 \to \rho^\pm \pi^\mp) = (27.6^{+8.4}_{-7.4} \pm 4.2) \times 10^{-6}.$$
(6)

These decays are sensitive to the form factors $F_0^{B\pi}$, $A_0^{B\rho}$, $A_0^{B\omega}$ and to the value of $N_c^{\rm eff}(LL)$. We consider two different form-factor models for heavy-to-light transitions: the BSW model ⁷ and the light-cone sum rule (LCSR) ⁸ and obtain $1.1 \leq N_c^{\rm eff}(LL) \leq 2.6$ from $\rho^0\pi^\pm$ and $\omega\pi^\pm$ modes. This is indeed what expected since the effective number of colors, $N_c^{\rm eff}(LL)$, inferred from the Cabibbo-allowed decays $B \to (D, D^*)(\pi, \rho)$ is in the vicinity of 2 (see ⁹) and since the energy released in the energetic two-body charmless B decays is

in general slightly larger than that in $B \to D\pi$ decays, it is thus anticipated that

$$|\chi(\text{two - body rare B decay})| \leq |\chi(B \to D\pi)|,$$
 (7)

and hence $N_c^{\text{eff}}(LL) \approx N_c^{\text{eff}}(B \to D\pi) \sim 2$.

Note that the branching ratio of $\rho^0\pi^\pm$ is sensitive to the change of the unitarity angle γ , while $\omega\pi^\pm$ is not. For example, we have $\mathcal{B}(B^\pm\to\rho^0\pi^\pm)\sim\mathcal{B}(B^\pm\to\omega\pi^\pm)$ for $\gamma\sim65^\circ$, and $\mathcal{B}(B^\pm\to\rho^0\pi^\pm)>\mathcal{B}(B^\pm\to\omega\pi^\pm)$ for $\gamma>90^\circ$. It appears that a unitarity angle γ larger than 90° , which is preferred by the previous measurement $^{10}\mathcal{B}(B^\pm\to\rho^0\pi^\pm)=(15\pm5\pm4)\times10^{-6}$, is no longer strongly favored by the new data of $\rho^0\pi^\pm$.

4 $B \rightarrow \pi\pi$ and πK Decays

The CLEO measurement of $B^0 \to \pi^+\pi^-$ mode [see Eq. (6)] puts a very stringent constraint on the form factor $F_0^{B\pi}$. Neglecting final-state interactions and employing $\gamma \equiv \operatorname{Arg}(V_{ub}^*) = 65^\circ$ and $|V_{ub}/V_{cb}| = 0.09$, and the effective number of colors $N_c^{\mathrm{eff}}(LL) = 2$, we find $F_0^{B\pi}(0) = 0.20 \pm 0.04$, which is rather small compared to the BSW value⁷ $F_0^{B\pi}(0) = 0.333$. This relatively small form factor will lead to two difficulties. First, the predicted $B \to K\pi$ branching ratios will be too small compared to the data as their decay rates are governed by the same form factor. Second, the predicted rate of $B \to K\eta'$ is also too small as the form factor $F_0^{BK}(0)$ cannot deviate too much from $F_0^{B\pi}(0)$, otherwise the SU(3)-symmetry relation $F_0^{BK} = F_0^{B\pi}$ will be badly broken.

There exist several possibilities that the $K\pi$ rates can be enhanced: (i) The CKM matrix element V_{ub} is small, say $|V_{ub}/V_{cb}| \approx 0.06$, so that the form factor $F_0^{B\pi}(0)$ is not suppressed ¹¹. However, this CKM matrix element $|V_{ub}/V_{cb}|$ is smaller than the recent LEP average ¹² $0.104^{+0.015}_{-0.018}$ and the CLEO result ¹³ $0.083^{+0.015}_{-0.016}$. (ii) A large nonzero isospin $\pi\pi$ phase shift difference of order 70° can yield a substantial suppression of the $\pi^+\pi^-$ mode³. However, a large $\pi\pi$ isospin phase difference seems to be very unlikely due to the large energy released in charmless B decays. Indeed, the Regge analysis of ¹⁵ indicates $\delta_{\pi\pi}=11^\circ$. (iii) Smaller quark masses, say $m_s(m_b)=65$ MeV, will make the (S-P)(S+P) penguin terms contributing sizably to the $K\pi$ modes but less significantly to $\pi^+\pi^-$ as the penguin effect on the latter is suppressed by the quark mixing angles. However, a rather smaller m_s is not consistent with recent lattice calculations. (iv) The unitarity angle γ larger than 90° will lead to a suppression of $B\to\pi^+\pi^-$ ^{14,3}, which in turn implies an enhancement of

 $F_0^{B\pi}$ and hence $K\pi$ rates. Therefore, the last possibility appears to be more plausible.

We find that the CLEO $K\pi$ and $\pi\pi$ data can be accommodated by $\gamma=105^\circ,^a_0F_0^{B\pi}(0)=0.28,~F_0^{BK}(0)=0.36,~N_c^{\rm eff}(LL)=2$. Note that $\gamma=(114^{+25}_{-21})^\circ$ is obtained by Hou, Smith, Würthwein ¹⁶ under the assumption of naive factorization. The calculated and experimental values of $K\pi$ decays are

$$\begin{split} \mathcal{B}(\overline{B}^0 \to K^- \pi^+) &= 18.6 \times 10^{-6}, & (17.2^{+2.5}_{-2.4} \pm 1.2) \times 10^{-6}, \\ \mathcal{B}(B^- \to \overline{K}^0 \pi^-) &= 17.0 \times 10^{-6}, & (18.2^{+4.6}_{-4.0} \pm 1.6) \times 10^{-6}, \\ \mathcal{B}(B^- \to K^- \pi^0) &= 12.6 \times 10^{-6}, & (11.6^{+3.0+1.4}_{-2.7-1.3}) \times 10^{-6}, \\ \mathcal{B}(\overline{B}^0 \to \overline{K}^0 \pi^0) &= 6.0 \times 10^{-6}, & (14.6^{+5.9+2.4}_{-5.1-3.3}) \times 10^{-6}. \end{split}$$

It is known that $K\pi$ modes are penguin dominated. As far as the QCD penguin contributions are concerned, it will be expected that $\mathcal{B}(\overline{B}^0 \to K^-\pi^+) \sim \mathcal{B}(B^- \to \overline{K}^0\pi^-)$ and $\mathcal{B}(B^- \to K^-\pi^0) \sim \mathcal{B}(\overline{B}^0 \to \overline{K}^0\pi^0) \sim \frac{1}{2}\mathcal{B}(B \to K\pi^\pm)$. However, as pointed out in 3,14 , the electroweak penguin diagram, which can be neglected in $\overline{K}^0\pi^-$ and $K^-\pi^+$, does play an essential role in the modes $K\pi^0$. With a moderate electroweak penguin contribution, the constructive (destructive) interference between electroweak and QCD penguins in $K^-\pi^0$ and $\overline{K}^0\pi^0$ renders the former greater than the latter; that is, $\mathcal{B}(B^- \to K^-\pi^0) > \frac{1}{2}\mathcal{B}(\overline{B}^0 \to \overline{K}^0\pi^-)$ and $\mathcal{B}(\overline{B}^0 \to \overline{K}^0\pi^0) < \frac{1}{2}\mathcal{B}(\overline{B}^0 \to K^-\pi^+)$ are anticipated. We see from Eq. (8) that, except for the decay $\overline{K}^0\pi^0$, the agreement of the calculated branching ratios for $K\pi$ modes with experiment is good. By contrast, the central value of $\mathcal{B}(\overline{B}^0 \to \overline{K}^0\pi^0)$ is much greater than the theoretical expectation. Since its experimental error is large, one has to await the experimental improvement to clarify the issue. The predicted pattern $K^-\pi^+ \gtrsim \overline{K}^0\pi^- \sim \frac{3}{2}K^-\pi^0 \sim 3\overline{K}^0\pi^0$ is consistent with experiment for the first three decays.

5 $B \to K\phi$, $K\eta'$ and $K^*\eta$ Decays

The decay amplitude of the penguin-dominated mode $B \to K\phi$ is governed by $[a_3 + a_4 + a_5 - \frac{1}{2}(a_7 + a_9 + a_{10})]$, where a_3 and a_5 are sensitive to $N_c^{\rm eff}(LL)$ and $N_c^{\rm eff}(LR)$, respectively. The current limit $\mathcal{B}(B^\pm \to \phi K^\pm) < 0.59 \times 10^{-5}$

^aThus far, only the data of $B \to \pi\pi$, $K\pi$ imply the possibility that $\cos \gamma < 0$. As noted in Sec. III, $\gamma > 90^{\circ}$ is not strongly favored by the measurements of $B \to \rho^0 \pi^{\pm}$, $\omega \pi^{\pm}$.

at 90% C.L. ¹⁰ implies that

$$N_c^{\text{eff}}(LR) \ge \begin{cases} 5.0 & \text{BSW}, \\ 4.2 & \text{LCSR}, \end{cases}$$
 (9)

with $N_c^{\rm eff}(LL)$ being fixed at the value of 2. Hence, we can conclude that $N_c^{\rm eff}(LR)>3>N_c^{\rm eff}(LL)$.

The improved measurements of the decays $B \to \eta' K$ by CLEO yield ⁶

$$\mathcal{B}(B^{\pm} \to \eta' K^{\pm}) = \left(80^{+10}_{-9} \pm 7\right) \times 10^{-6},$$

$$\mathcal{B}(B^{0} \to \eta' K^{0}) = \left(89^{+18}_{-16} \pm 9\right) \times 10^{-6}.$$
 (10)

This year CLEO has also reported the new measurement of $B\to K^*\eta$ with the branching ratios 17

$$\mathcal{B}(B^{\pm} \to \eta K^{*\pm}) = \left(26.4^{+9.6}_{-8.2} \pm 3.3\right) \times 10^{-6},$$

$$\mathcal{B}(B^{0} \to \eta K^{*0}) = \left(13.8^{+5.5}_{-4.6} \pm 1.6\right) \times 10^{-6}.$$
(11)

Theoretically, the branching ratios of $K\eta'$ ($K^*\eta$) are anticipated to be much greater than $K\pi$ ($K^*\eta'$) modes owing to the presence of constructive interference between two penguin amplitudes arising from non-strange and strange quarks of the η' or η . In general, the decay rates of $K\eta'$ increase slowly with $N_c^{\rm eff}(LR)$ if $N_c^{\rm eff}(LL)$ is treated to be the same as $N_c^{\rm eff}(LR)$, but fast enough with $N_c^{\rm eff}(LR)$ if $N_c^{\rm eff}(LL)$ is fixed at the value of 2. Evidently, the data much favor the latter case. As stressed in 3 , the contribution from the η' charm content will make the theoretical prediction even worse at the small values of $1/N_c^{\rm eff}$ if $N_c^{\rm eff}(LL) = N_c^{\rm eff}(LR)$! On the contrary, if $N_c^{\rm eff}(LL) \approx 2$, the $c\bar{c}$ admixture in the η' will always lead to a constructive interference irrespective of the value of $N_c^{\rm eff}(LR)$. The branching ratios of $K^*\eta$ in general decrease with $N_c^{\rm eff}(LR)$ when $N_c^{\rm eff}(LR)$ when $N_c^{\rm eff}(LR)$ but increase with $N_c^{\rm eff}(LR)$ when $N_c^{\rm eff}(LL) = 2$. Again, the latter is preferred by experiment. Hence, the data of both $K\eta'$ and $K^*\eta$ provide another strong support for a small $N_c^{\rm eff}(LL)$ and for the relation $N_c^{\rm eff}(LR) > N_c^{\rm eff}(LL)$. In other words, the nonfactorized effects due to (V-A)(V-A) and (V-A)(V+A) operators should be treated differently.

Several new mechanisms have been proposed in the past few years to explain the observed enormously large rate of $K\eta'$, for example, the large charm content of the η' ¹⁸ or the two-gluon fusion mechanism via the anomaly coupling of the η' with two gluons ^{19,20}. These mechanisms will in general predict a large rate for $K^*\eta'$ comparable to or even greater than $K\eta'$ and a very small rate for $K^*\eta$ and $K\eta$. The fact that the $K^*\eta$ modes are observed with sizeable branching ratios indicates that it is the constructive interference

of two comparable penguin amplitudes rather than the mechanism specific to the η' that accounts for the bulk of $B \to \eta' K$ and ηK^* branching ratios.

Finally, we would like to make a remark. As shown in 3 , the charged $\eta'K^-$ mode gets enhanced when $\cos\gamma$ becomes negative while the neutral $\eta'K^0$ mode remains steady. Therefore, it is important to see if the disparity between $\eta'K^\pm$ and $\eta'K^0$ is confirmed when experimental errors are improved and refined in the future.

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